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Inelastic J/Ψ Photoproduction off Nuclei: Gluon Enhancement or Double Color Exchange?

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Abstract

The nuclear enhancement observed in inelastic photoproduction of J/Ψ should not be interpreted as evidence for an increased gluon density in nuclei. The nuclear suppression of the production rate due to initial and final state interactions is calculated and a novel two-step color exchange process is proposed, which is able to explain the data.

1. Introduction

Measuring inelastic photoproduction of J/Ψ off heavy nuclei, $\gamma A \to J/\Psi X$, the EMC collaboration [1] and later the NMC collaboration [2] have observed an enhancement of the production rate compared to a proton or a light nuclear target: the experimental ratio of the cross sections per nucleon on tin to carbon exceeds the value one by about 10%. If this enhancement is fully attributed to an enhancement of the gluon density – as the authors of experiment [2] do – the gluon density in heavy nuclei should be larger by a similar value of a few percent, which is perfectly compatible with the theoretical expectations.

This conclusion is dangerous, however, since it neglects initial and final state interactions due to the presence of the other nucleons. The data from the E772 experiment [3] for J/Ψ production in proton-nucleus collisions show a suppression, which may be parameterized by $A^{-\epsilon}$, where $\epsilon = 0.08$. There is no reason, why this kind of initial/final state interaction should not be present in photoproduction. Therefore, the observed enhancement in the Sn/C ratio in $\gamma \to J/\Psi$ production of the order of 10% has to be seen in the light of an expected 20% suppression due to initial/final state interaction. It is theoretically unacceptable to claim that the discrepancy of 30% is attributed to a modified gluon distribution. Other mechanisms have to be looked for. This is the aim of the present paper, which is organized as follows.

In the next section the triple-Regge approach to inelastic photoproduction of charmonium is developed. Since the charmonium is produced diffractively via Pomeron exchange, we conclude that at $x_1 \to 1$ the reaction is sensitive to the double-gluon distribution function (similarly as in elastic photoproduction), rather than to the single gluon density [4]. Then, in section 3, the suppression factor for inelastic photoproduction of J/Ψ off nuclei is calculated using Glauber's eikonal approximation for the final state interaction. The formulas are derived for the first time. Since Glauber's approximation assumes lack of formation time we consider the color-singlet model for J/Ψ photoproduction in section 4. We demonstrate that contrary to the wide spread opinion the time interval for evaporation of color is quite short, about 1 fm in the energy range of the NMC experiment. Therefore Glauber's model

can be used in this case as well. Inelastic shadowing corrections to Glauber's approximation are considered in section 5. We demonstrate that this corrections make the nuclear medium substantially more transparent for charmonium: the effective absorption cross section turns out to be only about 60 % of that for J/Ψ. In section 6 we present the novel mechanism for diffractive photoproduction of charmonium off nuclei: the photon interacts with a bound nucleon via color-exchange and produces a $c\bar{c}$ pair in a color-octet state. The pair propagates through nuclear matter and looses energy via hadronization. Then, with another nucleon the $c\bar{c}$ pair exchanges color and converts to a color-singlet state. Such a process leads to J/Ψ's with $x_1 < 1$. Since this mechanism relies upon two interactions at least, its cross section may have an A^α-dependence with $\alpha > 1$. Interplay of this mechanism with the standard multiple scattering one nicely explains the NMC data.

2. Glauber theory for inelastic J/Ψ photoproduction off nuclei

The inelastic photoproduction cross section of a J/Ψ on a nucleus can be represented in the form

$$\frac{d\sigma_\Psi^{\gamma A}}{dx_1} = S_\Psi^{\gamma A}\ A\ \frac{d\sigma_\Psi^{\gamma N}}{dx_1}\ . \tag{1}$$

In this section we derive the nuclear suppression factor $S_\Psi^{\gamma A}$ within Glauber theory. There are two types of amplitudes which contribute to the inelastic photoproduction of a J/Ψ off nuclei:

I) The incoming photon does not interact in the nucleus until the point with coordinates (b, z) (the impact parameter and the longitudinal coordinate, respectively), where it interacts inelastically, $\gamma N \to J/\Psi X$, and produces the J/Ψ with the detected final momentum, which leaves the nucleus without inelastic interactions;

II) prior the inelastic interaction at the point (b, z) the photon produces elastically a J/Ψ at the point (b, z_1), $\gamma N \to J/\Psi N$, with $z_1 < z$. The J/Ψ propagates through the nucleus without interaction up to the point (b, z), where it produces the final J/Ψ in the reaction $J/\Psi N \to J/\Psi X$.

In both cases the longitudinal momentum transfer at the point (b, z) is so large, $q_{in} \approx (1 - x_1)\nu$, that only those amplitudes interfere, which correspond to the inelastic production of the J/Ψ with the desired momentum on the same nucleon.

On the other hand, the longitudinal momentum transfer in elastic photoproduction at the point (b, z_1) may be small at high energies,

$$q_{el} = \frac{Q^2 + M_\Psi^2}{2\nu} \ . \tag{2}$$

Therefore the amplitudes corresponding to different coordinates z_1 may interfere.

According to our classification of the amplitudes I and II the suppression factor $S_\Psi^{\gamma A}$ has three terms, $S_\Psi^{\gamma A} = (S_\Psi^{\gamma A})_1 + (S_\Psi^{\gamma A})_{12} + (S_\Psi^{\gamma A})_2$, where the first and the third terms correspond to squares of the amplitudes I and II, respectively, and the second one is the interference term.

The amplitude I squared provides the following suppression factor

$$\left(S_\Psi^{\gamma A}\right)_1 = \frac{1}{A} \int d^2b \int\limits_{-\infty}^{\infty} dz\ \rho(b, z)\ \exp\left[-\sigma_{in}^{\Psi N}\ T_z(b)\right] \ . \tag{3}$$

Here $T_z(b) = \int_z^\infty dz' \rho(b, z')$, where $\rho(b, z)$ is the nuclear density normalized to A. The nuclear thickness function is $T(b) = T_{z \to -\infty}(b)$.

The interference term has a form

$$\begin{aligned} \left(S_\Psi^{\gamma A}\right)_{12} &= -2\ \left(\frac{1}{2}\sigma_{tot}^{\Psi N}\right)\ \frac{1}{A} \int d^2b \int\limits_{-\infty}^{\infty} dz \rho(b, z)\ \exp\left[-\frac{1}{2}\left(\sigma_{in}^{\Psi N} - \sigma_{el}^{\Psi N}\right)\ T_z(b)\right] \\ &\times \int\limits_{-\infty}^{z} dz_1\ \rho(b, z_1)\ \cos[q_{el}(z - z_1)]\ \exp\left[-\frac{1}{2}\sigma_{tot}^{\Psi N}\ T_{z_1}(b)\right] \ . \end{aligned} \tag{4}$$

This term has a negative sign because the amplitude II plays the role of an absorptive correction to amplitude I.

To calculate the contribution of the amplitude II squared we can use the results of ref. [10, 11] for quasielastic photoproduction of vector mesons on nuclei, which we should apply to production of the J/Ψ prior the point (b, z). However, as different from [10, 11], we

should include both coherent (small q_T) and incoherent (large q_T) interactions of the J/Ψ in the nucleus. The result reads,

$$\begin{aligned}
\left(S_\Psi^{\gamma A}\right)_2 &= \frac{1}{A}\int d^2b \int\limits_{-\infty}^{\infty} dz\ \rho(b,z)\ \Biggl\{\sigma_{el}^{\Psi N} \int\limits_{-\infty}^{z} dz_1\ \rho(b,z_1)\ \exp\left[-\sigma_{in}^{\Psi N} T_{z_1}(b)\right] \\
&+ \frac{1}{2A}\ \sigma_{tot}^{\Psi N}\ (\sigma_{in}^{\Psi N} - \sigma_{el}^{\Psi N}) \int\limits_{-\infty}^{z} dz_1\ \rho(b,z_1)\ \exp\left[-\frac{1}{2}\sigma_{tot}^{\Psi N}\ T_{z_1}(b)\right] \\
&\times \int\limits_{z_1}^{z} dz_2\ \rho(b,z_2)\ \cos\left[q_{el}(z_2 - z_1)\right]\ \exp\left[-\frac{1}{2}\left(\sigma_{in}^{\Psi N} - \sigma_{el}^{\Psi N}\right)\ T_{z_2}(b)\right]\Biggr\}\ . \qquad (5)
\end{aligned}$$

At low energies, when $q_{el} \gg 1/R_A$ only $\left(S_\Psi^{\gamma A}\right)_1$ and the first term in curly brackets in (5) survive and lead to the nuclear suppression factor

$$S_\Psi^{\gamma A}|_{q_{el}\gg 1/R_A} = \frac{1}{\sigma_{in}^{\Psi N}}\ \frac{1}{A}\ \int d^2b\ \left\{\frac{\sigma_{tot}^{\Psi N}}{\sigma_{in}^{\Psi N}}\left[1 - e^{-\sigma_{in}^{\Psi N} T(b)}\right] - \sigma_{el}^{\Psi N}\ T(b)\ e^{-\sigma_{in}^{\Psi N} T(b)}\right\}\ . \qquad (6)$$

At high energies, in the limit $q_{el} \to 0$ the sum of expressions (4)-(6) simplifies very much,

$$S_\Psi^{\gamma A}|_{q_{el}\ll 1/R_A} = \frac{1}{A}\int d^2b\ T(b) e^{-\sigma_{in}^{\Psi N} T(b)}\ . \qquad (7)$$

Using the smallness of $\sigma_{tot}^{\Psi N} \approx 5\ mb$ we can expand (6) as $S_\Psi^{\gamma A}|_{q_{el}\ll 1/R_A} \approx 1 - \frac{1}{2}\sigma_{tot}^{\Psi N}\ \langle T\rangle$, where the mean nuclear thickness $\langle T\rangle = A^{-1}\int d^2b\ T^2(b)$. In the approximation eq. (7) we have kept only the first order of the expansion parameter $\sigma_{tot}^{\Psi N}\langle T\rangle$, and neglected $\sigma_{el}^{\Psi N}\langle T\rangle$, which is of the second order. To the same order the low-energy limit eq. (6) leads to $S_\Psi^{\gamma A}|_{q_{el}\gg 1/R_A} \approx 1 - \sigma_{tot}^{\Psi N}\ \langle T\rangle$.

Thus, the nuclear suppression factor at high energy turns out to be smaller than that at low energy. This result has a natural space-time interpretation: at high energies a photon may interact strongly via its hadronic fluctuations. The lifetime of the $c\bar{c}$ fluctuation of the photon, called coherence time, is $t_c \sim 1/q_{el}$. A fluctuation created long in advance propagates and attenuates through the whole nuclear thickness, as is explicitly exposed in (7). On the other hand, at low energy the photon converts to a $c\bar{c}$ inside the nucleus just prior to the J/Ψ production. So the $c\bar{c}$ pair propagates only through about a half

of the nuclear thickness. Such a contribution corresponds to $(S_\Psi^{\gamma A})_1$ given by (3). There is, however, a possibility to produce J/Ψ "elastically" in $\gamma N \to J/\Psi N$ prior the inelastic rescattering $J/\Psi N \to J/\Psi X$. This correction is suppressed by factor $\sigma_{el}^{\Psi N}/\sigma_{in}^{\Psi N}$ and is also included in eq. (6).

Expressions (3) - (5) give the correct extrapolation between the low-energy eq. (6) and the high-energy eq. (7) limits and are quite complicated. If one uses the same approximation $\sigma_{tot}^{\Psi N}\ \langle T\rangle \ll 1$ and expands eqs. (3) - (5) in this small parameter up to first order (compare with quasielastic photoproduction of J/Ψ off nuclei [12, 13]), one finds

$$S_\Psi^{\gamma A} \approx 1 - \frac{1}{2}\sigma_{tot}^{\Psi N}\ \langle T\rangle\ \Big[1 + F_A^2(q_{el})\Big]\ , \tag{8}$$

where the effective longitudinal nuclear formfactor squared is defined as

$$F_A^2(q_{el}) = \frac{1}{A\ \langle T\rangle}\ \int d^2b\ \left|\int\limits_{-\infty}^{\infty} dz\ \rho(b,z)\ \exp(iq_{el}z)\right|^2\ . \tag{9}$$

One can approximate $F_A^2(q_{el})$ quite well using a Gaussian parameterization instead of the more realistic Woods-Saxon form for the nuclear density [14] and obtains $F_A^2(q_{el}) = \exp(-R_A^2\ q_{el}^2)/3$, where R_A^2 is the mean squared charge radius of the nucleus [14]. Note that $S_\Psi^{\gamma A}$ is independent of x_1, therefore this mechanism of nuclear suppression keeps the same x_1-distribution of the cross section as on a free nucleon and cannot account for the data (cf. section 4).

There is another possibility within the standard Glauber approach to produce the J/Ψ in two-step process on a nucleus: the photon produces the J/Ψ inelastically on one nucleon with a subsequent diffractive scattering of the charmonium on another nucleon. Inelasticities of these two interactions should be adjusted to provide the desired x_1 for the final J/Ψ. The contribution of this mechanism to the cross section is evaluated in the Appendix and a value of less than 2% is found, which we neglect compared to the main mechanism considered above.

3. Triple-Regge formalism for J/Ψ photoproduction.

As soon as the nuclear suppression factor $S_{\Psi}^{\gamma A}$ is calculated within Glauber approximation, one can predict the nuclear cross section provided that the J/Ψ photoproduction cross section on a free nucleon is known. We review briefly in this and the next sections the well accepted mechanisms in order to justify the usage of Glauber model and to get a reasonable parameterization for the cross section.

In the standard triple-Regge formalism (see for instance [5]) the diffractive inelastic photoproduction of J/Ψ is described by the triple-Pomeron graph, depicted in Fig. 1.

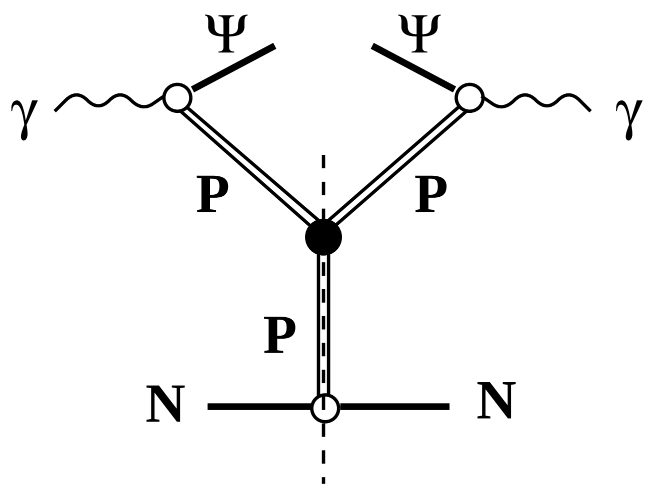

Figure 1: The triple-Pomeron graph for the cross section of the reaction $\gamma N \to J/\Psi X$. The dashed line shows that only the absorptive part of the amplitude is included.

The corresponding differential cross section reads

$$\frac{d\sigma(\gamma N \to J/\Psi X)}{dx_1\ dt} = G(\gamma N \to J/\Psi X)\frac{\exp(-B_{in}^{\Psi p} p_T^2)}{(1-x_1)^{2\alpha_P(t)-1}}\ . \tag{10}$$

Here x_1 and t are the relative share of the photon light-cone momentum carried by the J/Ψ and the 4-momentum squared of the produced J/Ψ, respectively. $\alpha_P(t)$ is the Pomeron trajectory and $B_{in}^{\Psi p}$ is the slope parameter for inelastic $J/\Psi - p$ collisions. The factor $G(\gamma N \to J/\Psi X)$ includes the triple-Pomeron coupling as well as the $P\gamma\Psi$ and the PNN vertices. It can be evaluated using the data from hadronic interactions, assuming factorization of the Pomeron.

$$G(\gamma N \to J/\Psi X) = G(pp \to pX) \; \frac{\sigma(\gamma p \to J/\Psi p) \; B_{el}^{\Psi p}}{\sigma_{el}^{pp} \; B_{el}^{pp}} \; . \qquad (11)$$

The momentum transfer squared t is rather large in $\gamma N \to J/\Psi X$, $|t|_{min} = (M_\Psi^2 + x_1 Q^2)(1-x_1)/x_1$, where $x_1 = p_\Psi^+/p_\gamma^+$ ($\approx x_F$ at $x_1 \to 1$) is the fraction of the photon light-cone momentum carried by the J/Ψ. $|t_{min}|$ may reach a few GeV^2 dependent on x_1. The Pomeron trajectory $\alpha_P(t)$ is known to approach the value of 1 at high $|t|$ [6], therefore we take $\alpha_P(t) = 1$.

Combining all the factors we get an estimate for the x_1-distribution of the inclusive photoproduction of J/Ψ on a proton, integrated over transverse momenta

$$\frac{d\sigma(\gamma N \to J/\Psi X)}{dx_1} = \frac{B_{el}^{\Psi p}}{B_{el}^{pp} \; B_{in}^{\Psi p}} \; \frac{\sigma(\gamma p \to J/\Psi p)}{\sigma_{el}^{pp}} \; \frac{G(pp \to pX)}{1 - x_1} \; . \qquad (12)$$

The effective triple-Pomeron constant $G(pp \to pX) = 3.2 \; mb/GeV$ is fixed by an analysis of data on $pp \to pX$ [5]. The elastic pp cross section and the slope parameter are taken as $\sigma_{el}^{pp} = 8 \; mb$ and $B_{el}^{pp} \approx 10 \; GeV^{-2}$. The slope of the p_T-distribution of elastic J/Ψ photoproduction is assumed $B_{el}^{\Psi N} \approx B_{el}^{pp}/2$, since the slopes are related to the nucleon and the J/Ψ form factors respectively. On the other hand, in the inelastic photoproduction of J/Ψ the nucleon is destroyed and its formfactor does not contribute to the effective slope (the same as in the case of $pp \to pX$ reaction, where the slope $B_{in}^{pp} \approx 4 \; GeV^{-2}$ [5]). Therefore $B_{in}^{\Psi N}$ is only due to the charmonium formfactor. It was measured in the EMC experiment [7] to be $B_{in}^{\Psi p} = 0.58 \pm 0.07 \; GeV^{-2}$, but the NMC experiment [8] led to $B_{in}^{\Psi p} = 0.23 \pm 0.02 \; GeV^{-2}$. For a mean photon energy $\langle \nu \rangle \approx 100 \; GeV$, the elastic photoproduction cross section is $\sigma(\gamma p \to J/\Psi p) \approx 12 \; nb$ according to the EMC measurements [7], but is $\approx 27 \; nb$ according to the NMC results [9]. In view of the uncertainties in the measured values of the elastic photoproduction cross section and the inelastic slope we fix them at a middle value $\sigma(\gamma p \to J/\Psi p) = 20 \; nb$ and $B_{in}^{\Psi p} = 0.45 \; GeV^{-2}$.

The the cross section formula (12) is calculated with the parameters given above and is compared with the data from the NMC experiment [9] in Fig. 2. A good agreement is found both in shape and normalization.

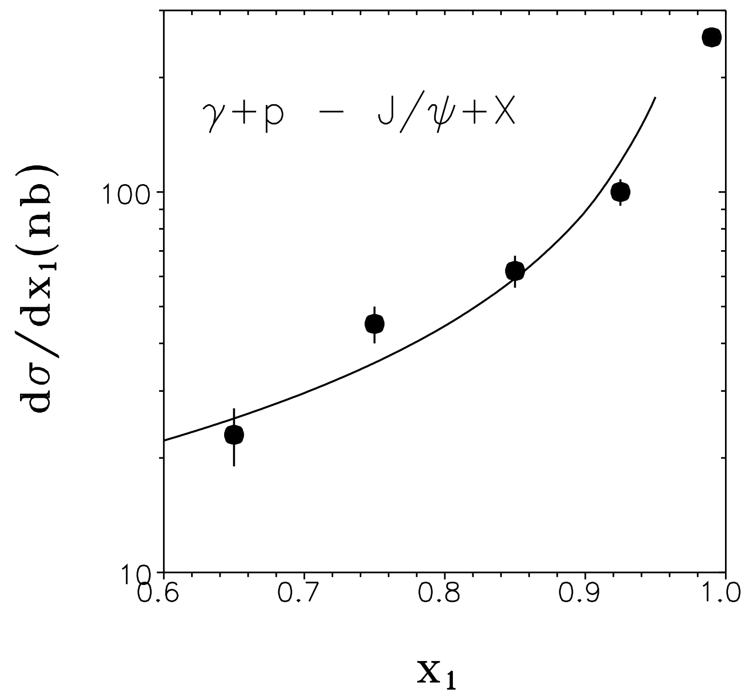

Figure 2: The data from the NMC experiment [2] for inelastic J/Ψ photoproduction on protons. The solid curve shows predictions based on the triple-Pomeron phenomenology.

Note that the triple-Pomeron is known to dominate in the inclusive cross section of pp interaction only at $x_1 > 0.9$ [5]. The nondiffractive RRP triple-Reggeon graphs are important at smaller x_1. OZI rule suppresses the RRP term for J/Ψ interactions, therefore the triple Pomeron contribution may be important down to lower x_1. Nevertheless, the triple-Pomeron phenomenology does not necessarily describe the data for $x_1 < 0.8$. Particularly, the longitudinal formfactor of the $\gamma \to J/\Psi$ vertex may substantially diminish the cross section at mid x_1. Yet we need a parameterization of the data for further applications, and Fig. 2 shows that eq. (12) reproduces the data well down to $x_1 = 0.6$.

Usually the inelastic photoproduction of J/Ψ is treated perturbatively, within the so called color-singlet model [4]. In that approach the cross section is predicted to be proportional to the low-x gluon density distribution in the proton. Although the Pomeron by definition includes all possible kinds of perturbative graphs, but the ones of the color-singlet

model are not the leading contributions at $x_1 \to 1$. Therefore the spectrum of J/Ψ predicted by the color singlet model does not peak at $x_1 \to 1$, while the triple-Pomeron contribution does. The dominant mechanism comes from the ladder-type graphs with excitation of the proton. Correspondingly, the cross section is proportional to the double-gluon density, which is approximately equal to the single-gluon density squared.

We conclude that for $x_1 \to 1$ the inelastic photoproduction of J/Ψ cannot be used as a probe for the single-gluon distribution, but it might be possible at mid x_1-values.

4. Color-singlet model

A popular approach to inelastic diffractive photoproduction of J/Ψ on a nucleon target is the color-singlet model (CSM) [4]. It is somewhat similar to the mechanism suggested in section 6, except that the color-octet $c\bar{c}$ pair becomes colorless via radiation of a gluon, rather than by means of a color-exchange rescattering.

Radiation of a gluon takes some time, which may be important if the process takes place in a nuclear environment. Indeed, rescatterings of the color-octet $c\bar{c}$ pair may induce additional energy loss, and may result in an effective A-dependence. Therefore we need to understand the space-time pattern of gluon radiation in the CSM.

Sometimes the production of a charmonium in CSM is associated with a two-step space-time development: first the color-octet $c\bar{c}$ pair is produced, then it propagates and finally converts into a color-singlet state via radiation of a gluon. The second stage is usually assumed to take a long time.

We do not think that such a treatment is possible in view of a principal quantum-mechanical uncertainty: analogous to the standard electromagnetic bremsstrahlung one cannot say whether the gluon/photon has been radiated before or after the interaction. Both Feynman graphs contribute to the cross section (actually both are taken into account in perturbative calculations [4]). One can understand this uncertainty as related to the lifetime of a fluctuation on the initial particle containing the radiated gluon and the color-octet $c\bar{c}$ pair. This lifetime given by the energy-denominator equals to

$$t_r = \frac{2\nu}{M^2} \ , \tag{13}$$

where M is the effective mass of the fluctuation,

$$M^2 = \frac{M_{c\bar{c}}^2}{x_1} + \frac{k_T^2}{x_1(1-x_1)} \ . \tag{14}$$

Taking the mass of the $c\bar{c}$ pair $M_{c\bar{c}} \approx M_\Psi$ we get the following expression for the radiation time

$$t_r = \frac{2\ \nu\ x_1(1-x_1)}{k_T^2 + M_\Psi^2(1-x_1)} \ . \tag{15}$$

This radiation time is only about 1 fm in the kinematical region of the NMC experiment. One can safely neglect an influence of nuclear medium on the process of hadronization for such a short time interval. Therefore the multiple scattering approach developed in sections 2 can be applied in the case of CSM.

Note that even at very high energies, when t_r becomes long, the effects of induced radiation by the color-octet $c\bar{c}$ pair during radiation time can be neglected since $2\kappa_{ind}/M_\Psi^2 \ll 1$, where $\kappa_{ind} \leq 1\ GeV^2$ is the density of energy loss via induced radiation.

5. Beyond Glauber approximation

Considering photoproduction of J/Ψ in the quark representation one understands that a colorless $c\bar{c}$ wave packet, rather than a J/Ψ, is produced and propagates through the nucleus, but the J/Ψ may be formed far away from the nucleus if the energy is high. Due to quark-hadron duality one can consider the same effects taking into account photoproduction of higher diffractive excitations of the charmonium, Ψ', Ψ'' etc., as well as all diagonal and off-diagonal diffractive rescatterings of these states in nuclear matter. The corresponding corrections are know as inelastic shadowing [15]. The simplest, but quite accurate two-coupled-channel approximation including J/Ψ and Ψ' was considered in [16]. It was found that the the $c\bar{c}$ wave packet produced in pp interaction is such a combination of J/Ψ and Ψ', which is quite close to the eigenstate of interaction having the minimal possible interaction

cross section. This fact explains particularly why J/Ψ and Ψ experience similar nuclear attenuation (cf. the E772 experiment [3]).

We assume that the ratio of the photoproduction amplitudes

$$R_{\gamma p} = \frac{\langle\Psi'|\widehat{f}|\gamma\rangle}{\langle\Psi|\widehat{f}|\gamma\rangle} \ , \tag{16}$$

where $\widehat{f}$ is the operator of diffractive amplitude, is close to one. In pp scattering the corresponding value of R_{pp} is measured to be 0.5 (see in [16]). This assumption on $R_{\gamma p}$ can be motivated by the similarity of the $c\bar{c}$ components of the wave functions of a photon and of a gluon.

The solution of evolution equation for the $c\bar{c}$ wave packet propagated through nuclear matter, which incorporates with both effects of the coherence and formation lengths will be presented elsewhere [18], but some numerical results are published in [17]. Here we use the approximation of small $\sigma_{tot}^{\Psi N}\langle T\rangle \ll 1$. Then one can modify the Glauber model expression (8) as follows [17]

$$S_{\Psi}^{\gamma A} \approx 1 - \sigma_{eff}T_{eff} \approx exp(-\sigma_{eff}T_{eff}) \ , \tag{17}$$

where

$$T_{eff} = \frac{1}{2}\langle T\rangle[1 + F_A^2(q_c)] \tag{18}$$

is according to eq. (8) the effective nuclear thickness covered by the charmonium, which is controlled by the coherence length $l_c \approx 1/q_{el}$ [11], and

$$\sigma_{eff} = \sigma_{in}^{\Psi}[1 + \epsilon\ R\ F_A^2(q_f)] \tag{19}$$

is the effective absorptive cross section for charmonium, which is controlled by the formation length for the charmonium wave function, $l_f \approx 1/q_f$, where

$$q_f = \frac{M_{\Psi'}^2 - M_{\Psi}^2}{2x_1\nu}, \tag{20}$$

is the longitudinal momentum transfer in the off-diagonal reaction $\Psi N \rightleftharpoons \Psi' N$.

The new parameter ϵ in eq. (19) is

$$\epsilon = \frac{\langle \Psi' | \widehat{f} | \Psi \rangle}{\langle \Psi | \widehat{f} | \Psi \rangle} \ , \tag{21}$$

and is estimated in [16] in the oscillator approximation, $\epsilon = -\sqrt{2/3}$. Remarkably, with these parameters at high energies $\sigma_{eff} \approx 0.6 \ \sigma_{in}^{\Psi N}$, i.e. the J/Ψ attenuates in nuclear matter substantially less than predicts Glauber model. This may be understood as a natural consequence of color transparency.

We calculate ratio of production rates of J/Ψ for tin to carbon using eq. (17). The results for $\nu = 70$ and $100 \ GeV$ are plotted in Fig. 3 in comparison with the data from the NMC experiment [2]. We see that the discrepancy with the inelastic photoproduction data ($x_1 < 0.85$) is so large that it is difficult to explain it by a few percent enhancement of the gluon density in heavy nuclei. One has to look for another explanation of the EMC-NMC effect.

6. Novel mechanism of J/Ψ photoproduction in nuclear matter

According to the arguments in section 3 the photoproduction of J/Ψ on a nucleon proceeds via Pomeron exchange, which in the language of perturbative QCD can be interpreted as a colorless gluonic exchange, containing at least two gluons in the t-channel. In a nuclear interaction this colorless exchange can be split into two color-exchange interactions with different nucleons, keeping the beam and the nuclear remnants in the final state colorless [21, 22].

The same physics can be represented as a consequence of unitarity: the imaginary part of the elastic forward scattering amplitude of J/Ψ on a nucleon is a product of an inelastic $J/\Psi \ N$ amplitude with itself time-conjugated (up to small elastic and diffractive corrections). In the case of a nuclear target one has to consider also products of two inelastic amplitudes on different nucleons.

Thus, the J/Ψ photoproduction on a nucleus can result from the photoproduction on a bound nucleon of a $c\bar{c}$ pair in a color-octet state, which propagates through nuclear matter

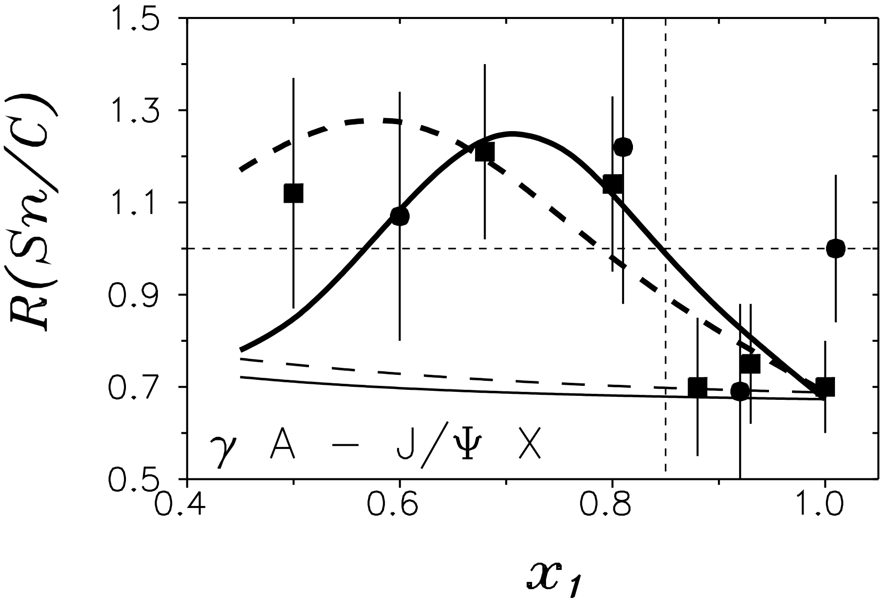

Figure 3: The data from the NMC experiment [2] for the ratio of inelastic J/Ψ photoproduction rates on tin to carbon, measured with 200 GeV (squares) and 280 GeV (circles) muons. The thin solid and dashed curves correspond to the Glauber-Gribov mechanism for nuclear suppression of J/Ψ and $\nu = 100$ and 70 GeV, respectively. The thick solid and dashed curves show predictions based on the proposed novel mechanism of diffractive interaction with a nucleus at $\nu = 100$ and 70 GeV, respectively. For the origin of the dividing line at $x_1 = 0.85$ see text.

loosing energy via hadronization until another color-exchange interaction inside the nucleus converts the pair back to a colorless state. Actually, this is a specific QCD mechanism of diffractive excitation of the nucleus. The energy of the J/Ψ (and the excitation energy of the nucleus) is less than the photon energy by the amount lost via hadronization by $c\bar{c}$ pair while it is in the color octet state. We assume a constant energy loss per unit of length $dE/dz = -\kappa$, what is true both for the color string [23] and for the bremsstrahlung [24] mechanisms of hadronization. In this case we have

$$x_1 = 1 - \frac{\kappa}{\nu}\Delta z \ , \tag{22}$$

where Δz is the longitudinal distance covered by the $c\bar{c}$ pair in the color-octet state.

The correction from this mechanism to the cross section for nuclear photoproduction can be calculated as follows

$$\begin{aligned} \Delta\left(\frac{d\sigma_{\Psi}^{\gamma A}}{dx_1}\right) &= 2\pi\,\frac{\nu}{\kappa}\,B_{el}^{\Psi N}\,\sigma(\gamma N \to J/\Psi N)\int d^2b \int\limits_{-\infty}^{\infty} dz_1\ \rho_A(b,z_1)\int\limits_{z_1}^{\infty} dz_2\ \rho_A(b,z_2) \\ &\delta\left[z_2 - z_1 - (1-x_1)\frac{\nu}{\kappa}\right]\ \exp\left\{-\sigma_{in}^{\Psi N}\left[T(b) - T_{z_1}(b) + T_{z_2}(b)\right]\right\}\ . \end{aligned} \tag{23}$$

The derivation of this expression includes following ingredients. We assume that the photon energy is sufficiently high, $\nu \gg M_{\Psi}^2 R_A/2$ to allow the photon to convert into the $c\bar{c}$ pair long in advance. At the first interaction point (b, z_1) the $c\bar{c}$ pair converts into the color octet state, corresponding to the cut Pomeron with cross section $\sigma_{in}^{\Psi N}$. It is important that no longitudinal momentum is transferred to the $c\bar{c}$ pair at this point. All the loss of longitudinal momentum, $(1 - x_1)\nu$, is due to the hadronization, until the last interaction at point (b, z_1), where the color octet $c\bar{c}$ pair converts back to the colorless state. This interaction proceeds with the same cross section as the very first one, except for a factor $1/8$ (no summation over final colors of the $c\bar{c}$ in this case). We also use the relation $\sigma_{el}^{\Psi N} = (\sigma_{tot}^{\Psi N})^2/16\pi B_{el}^{\Psi N}$, where $B_{el}^{\Psi N} \approx 5\ GeV^{-2}$ is the slope of ΨN elastic scattering (see section 3. The δ-function in eq. (23) takes care of energy-conservation including the energy loss (22) during the intermediate state hadronization. We also assume that the $c\bar{c}$ pair does not attenuate when it propagates through the nucleus in a color octet state. Actually, rescatterings of such a pair also may cause an attenuation [25], but it is small provided that the photon energy is sufficiently high $(1 - x_1)\nu/\kappa \gg R_A$. However, the $c\bar{c}$ pair does attenuate propagating through the nucleus in the colorless state before the first interaction (if $q_{el}R_A \ll 1$) and after the last one. We assume in eq. (23) a long coherence length and zero formation length. We neglect the energy variation due to these effects because the estimate eq. (23) is quite rough and the accuracy of data [1, 2] is low as well.

The only unknown parameter we are left with in eq. (23) is κ, the energy loss per unit length. In the string model it equals the string tension. The color-octet string tension κ may substantially exceed the known value for color-triplet strings $\kappa_3 \approx (2\pi\alpha'_R)^{-1} \approx 1\ GeV/fm$ [23], where $\alpha'_R \approx 1\ GeV^{-2}$ is the universal slope of Regge trajectories. Indeed, the string tension of the color-octet string must to be related to the slope of the Pomeron trajectory, $\alpha'_P \approx 0.2\ GeV^{-2}$, which is much smaller than α'_R. Thus, $\kappa_8 \approx (2\pi\alpha'_P)^{-1} \approx 5\ GeV/fm$ [22].

One may also treat the energy loss perturbatively as a result of gluon bremsstrahlung by the color-octet $c\bar{c}$ pair. It was demonstrated in [24] that bremsstrahlung provides a constant density of energy loss, similar to the string model. The mean squared color-octet charge is 9/4 times bigger than the color-triplet one, so is the energy loss. Besides, in the case under discussion the color-charge appears during the first interaction with a bound nucleon at the time $t = t_1$, then propagates for the time, $\Delta t = t_2 - t_1$ and disappears when the $c\bar{c}$ pair becomes colorless. Such a process with switch-on and switch-off causes a double bremsstrahlung [24] (see also [21]): only that part of the frequency spectrum is emitted during a time interval Δt, which has radiation time $t_r = 2\omega/k_T^2 < \Delta t$, where ω and k_T are the energy and transverse momentum of the radiated gluon, respectively. At the time $t = t_2$ when the $c\bar{c}$ pair converts to a colorless state and the radiation process stops. A new radiation process caused by the charge stopping starts, and another piece of the gluon spectrum is radiated, identical to the one previously emitted. Thus, the intensity of radiation and the energy loss are twice as large as in a single-scattering process. Summarizing, we expect the energy loss for gluon bremsstrahlung in the double-scattering process with a color-octet intermediate state to be 4.5 times larger than that for radiation of a color-triplet charge, produced in a single scattering process. This estimate is very close to the one we found above in the string model, so we fix for further calculations $\kappa = 5\ GeV/fm$.

The new mechanism provides an x_1-dependence of the cross section quite different from the standard one. There is no cross section for $x_1 < 1 - \kappa R_A/\nu$, since the amount of lost energy is restricted by the length of the path of the color-octet $c\bar{c}$ pair inside the nucleus. This contribution does not peak at $x_1 \to 1$, but may have even a minimum due to a longer

path of the colorless $c\bar{c}$ pair in nuclear matter.

The contribution of this mechanism may cause an A-dependence steeper than A^1, since there is more room for integration over longitudinal coordinates z_1 and z_2. For this reason the contribution of the mechanism under discussion, is important for heavy nuclei at moderately high x_1. We compare the contribution of Glauber-Gribov mechanism calculated with eq. (17) and the novel one given by eq. (23) to J/Ψ photoproduction on lead at different photon energies in Fig. 4. We see a strong energy variation of the latter contribution, which shifts with energy towards $x_1 = 1$. At the same time the energy dependence of nuclear effects provided by a variation of the coherence and formation lengths is so small that we plotted only one curve for the Glauber-Gribov mechanism at $\nu = 100\ GeV$.

We see that both mechanisms the Glauber one and the two-step production are of the same order at $x_1 \approx 0.6 - 0.7$.

The novel mechanism proposed here does not display the Feynman scaling typical for the standard mechanism, but this contribution grows with energy at fixed x_1. This is easy to understand: the total contribution integrated over x_1 is energy independent, but the range of $1 - x_1$ shrinks $\propto 1/\nu$.

With the above value of $\kappa = 5\ GeV/fm$ and $\sigma_{in}^{\Psi N} = 5\ mb$ we calculate the contribution of the double-color-exchange mechanism (23). Combining with the Glauber contribution from section 2 we have

$$\frac{d\sigma_{\Psi}^{\gamma A}}{dx_1} = S_{\Psi}^{\gamma N}\ \frac{d\sigma_{\Psi}^{\gamma N}}{dx_1} + \Delta\left(\frac{d\sigma_{\Psi}^{\gamma A}}{dx_1}\right)\ . \tag{24}$$

The result for the ratio of the cross sections for tin to carbon is plotted in Fig. 3 for two photon energies $\nu = 70$ and $100\ GeV$ typical for the NMC experiment in comparison with the data [2].

It should be noted that such a comparison is legitimate only for $x_1 < 0.85$ where the experimental resolution of the J/Ψ momentum suffices to separate events of inelastic photoproduction from the quasielastic ones. For $x_1 > 0.9$ quasielastic photoproduction of J/Ψ is supposed [1, 2] to be the dominant process. In this region we should compare the data

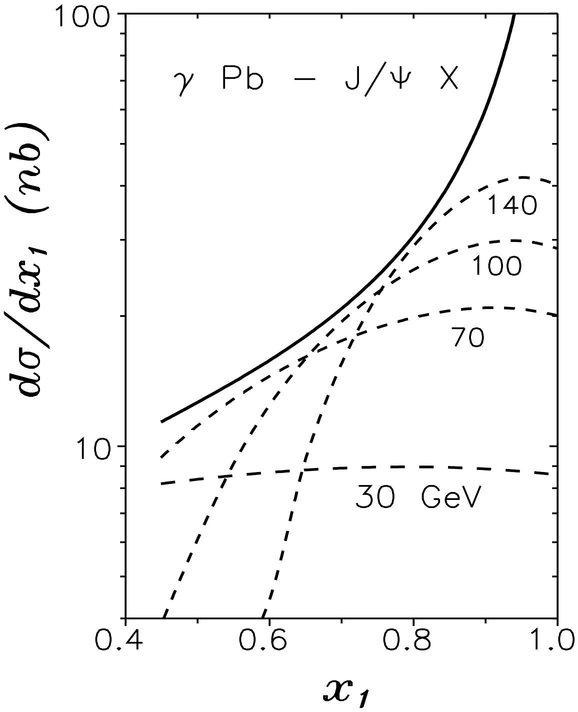

Figure 4: The cross section per nucleon for J/Ψ photoproduction on lead. The solid curve corresponds to Glauber-Gribov mechanism of nuclear shadowing at $\nu = 100\ GeV$. The novel mechanism contributions at $\nu = 70$, 100 and 140 GeV are shown by dashed curves.

with the quasielastic photoproduction cross section. The corresponding formula is derived in [10, 11]. It is different from formulas (3) - (5) for inelastic photoproduction, but in the approximation $\sigma_{tot}^{\Psi N}\langle T\rangle \ll 1$ looks analogous to eq. (17)

$$\sigma_{Qel}(\gamma A \to \Psi) \approx \sigma(\gamma N \to \Psi N)\ A\ \left\{1 - \frac{1}{2}\sigma_{tot}^{\Psi N}\ \langle T\rangle\ \left[1 + F_A^2(q_{el})\right]\ \left[1 + \epsilon\ R\ F_A^2(q_f)\right]\right\}\ , \tag{25}$$

The results of calculation with eq. (17) shown in Fig. 3 correspond to quasielastic scattering in the limit $x_1 \to 1$, where the momentum transfer eq. (20) acquires the value $q_f = (M_{\Psi'}^2 - M_{\Psi}^2)/2\nu$.

Fig. 3 shows a good agreement with the data both for inelastic and quasielastic photoproduction of J/Ψ off nuclei.

It is important to study the energy dependence of the effect under discussion, which is expected to be interesting. It seems most likely that new data on inelastic J/Ψ photoproduction will be taken with the HERMES spectrometer at HERA. In Fig. 5 we show predictions for nuclear suppression/enhancement factor $S_{\Psi}^{\gamma A}$ at $\nu = 30\ GeV$ for xenon and nitrogen.

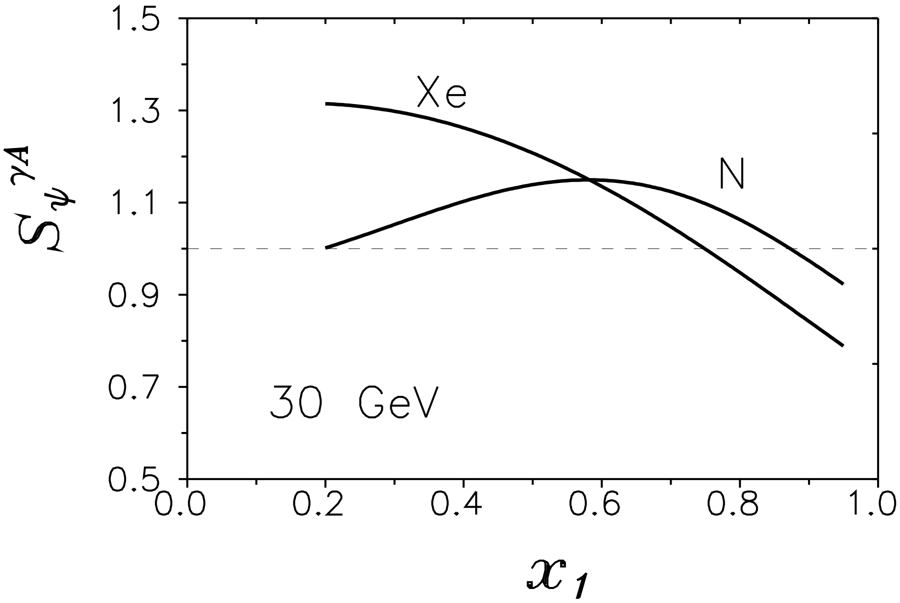

Figure 5: Nuclear suppression/enhancement factor for J/Ψ photoproduction on xenon and nitrogen at $\nu = 30\ GeV$

7. Conclusions and discussion

We summarize the main results of the paper as follows.

- The cross section of inelastic photoproduction of charmonia off nuclei is calculated within the eikonal (Glauber) approximation. Corresponding formulas are derived for the first time. We also correct the Glauber formulas for inelastic shadowing, which makes nuclear medium substantially more transparent for charmonium.

- The experimentally observed nuclear enhancement of inelastic J/Ψ photoproduction at $x_1 < 0.85$ is found to be too large compared with expectation of Glauber model to be explained by an enhancement of gluons in nuclei at low x_{Bj}.

- A novel QCD mechanism of charmonium photoproduction is suggested, which is not included in the standard multiple-scattering theory and can only occur in a nuclear target. This mechanism arises as a natural consequence of the colored structure of the Pomeron: the projectile hadron (or a hadronic fluctuation of the photon) induces a color dipole in the target nucleus, leaving the nucleus in a colorless state. This is a specific way of diffractive excitation of the nucleus and inclusive production of leading hadrons.

- As different from the Glauber model the new mechanism contribution is essentially energy-dependent and violates Feynman scaling. Due to steeper an A-dependence this mechanism successfully competes with Glauber one at moderate x_1 in the energy range of the EMC and NMC experiments and nicely explains the observed [1, 2] nuclear enhancement of the photoproduction cross section. This effect is expected to be squeezed towards $x_1 = 1$ and to vanish under the quasielastic peak at higher energies.

- We conclude that the observed nuclear enhancement of inelastic photoproduction of J/Ψ is the evidence of a new phenomenon related to specifics of QCD. An analogous effect in the case of interaction of light hadrons turns out to be suppressed and escapes observations [21, 22]. Photoproduction of heavy quarkonia might be a unique reaction where the new mechanism can be observed.

Note, that a specific channel of decay of the color-dipole created inside the nucleus into two nucleons was suggested in [26] as a source of the backward protons in the laboratory frame. Measurements by the SIGMA Collaboration [27] of backward proton production in pion-beryllium interaction at 40 GeV nicely confirmed the theoretical expectations. This is another manifestation of the of double-color-exchange mechanism.

The important signature of the mechanism under discussion its specific energy dependence. The nuclear enhancement does not vanish at higher energies, but is squeezed towards $x_1 = 1$. This is different from gluon distribution in a nucleus, which anyway have to switch

to nuclear shadowing at higher energies, i.e. smaller Bjorken x_2.

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Appendix

We estimate the correction $\delta(d\sigma_{\Psi}^{\gamma A}/dx_1)$ from inelastic production of J/Ψ with inelasticity α at one point in the nucleus, with subsequent inelastic rescattering with inelasticity β on another bound nucleon. This contribution to the cross section integrated over transverse momenta reads

$$\delta\left(\frac{d\sigma_{\Psi}^{\gamma A}}{dx_1}\right) \approx A\,\widetilde{T}\int\limits_0^1 d\alpha\int\limits_0^1 d\beta\;\delta(x_1-\alpha\beta)\,\frac{d\sigma_{\Psi}^{\gamma N}}{d\alpha}\,\frac{d\sigma_{\Psi}^{\gamma N}}{d\beta}\;, \tag{26}$$

where

$$\widetilde{T} = \frac{1}{2A}\int d^2b\; T^2(b)\exp[-\sigma_{in}^{\Psi N}T(b)] \tag{27}$$

We assume for the sake of simplicity that the energy is high, $q_{el}R_A \ll 1$.

After integration (26) takes the form

$$\delta\left(\frac{d\sigma_{\Psi}^{\gamma A}}{dx_1}\right) = \widetilde{T}\;K(x_1)\;A\;\frac{d\sigma_{\Psi}^{\gamma N}}{dx_1}\;, \tag{28}$$

noi where the factor $K(x_1)$ is a slow function of x_1,

$$K(x_1) = \left(\frac{\sigma_{tot}^{\Psi p}}{\sigma_{tot}^{pp}}\right)^2\;\frac{G(pp\to pX)}{B_{in}^{\Psi p}}\;\left[(1+x_1)\;\ln\left(\frac{1-x_1}{\epsilon}\right)+x_1\;ln\left(\frac{1}{x_1}\right)\right] \tag{29}$$

Here $\epsilon = M_0^2/s$, where M_0 is the minimal mass of the excited nucleon in final state, which we include in the integration over inelasticities in (26). At $M_0 = 2\;GeV$ and $x_1 = 0.9$ we estimate the factor $K \approx 0.03\;fm^2$

The effective nuclear thickness (27) for tin is $\widetilde{T}^{Sn} \approx 0.3\ fm^{-2}$.

We conclude that according to these estimates and expression (29) the correction under discussion has nearly the same form of x_1-dependence as the standard mechanism contribution, but is suppressed by the factor $K\ \widetilde{T}\ A/S_{\Psi}^{\gamma N} \approx 0.016$. We can neglect such a small correction for our purposes.

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